

Dual Higgs Mechanism and Nonperturbative QCD

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We study the dual Higgs mechanism for the nonperturbative QCD. We point out two kind of the “see-saw” relations between electric and magnetic sectors. Owing to the Dirac condition $eg = 4\pi$, the dual Ginzburg-Landau theory has the asymptotic freedom nature on the gauge coupling constant e , where the “walking coupling constant” is predicted for e in the infrared region. We study also QCD-monopoles and instantons, which are two relevant topological objects in QCD. Strong correlations between the QCD-monopole trajectory and the topological charge are found both in the continuum theory and in the lattice gauge theory.

1. Introduction

Quantum chromodynamics (QCD) is the fundamental theory of the strong interaction [1-3]. In spite of the simple form of the QCD lagrangian,

$$\mathcal{L}_{\text{QCD}} = -\frac{1}{2}\text{tr}G_{\mu\nu}G^{\mu\nu} + \bar{q}(\not{D} - m_q)q, \quad (1)$$

it miraculously provides quite various phenomena like color confinement, dynamical chiral-symmetry breaking, non-trivial topologies, quantum anomalies and so on (Fig.1). It would be interesting to compare QCD with the history of the Universe, because a quite simple ‘Big Bang’ also created various things including galaxies, stars, lives and thinking reeds. Therefore, QCD can be regarded as an interesting miniature of the history of the Universe. This is the most attractive point of the QCD physics.

Since it is quite difficult to understand the various QCD phenomena and their underlying mechanism at the same time, many methods and models have been proposed to understand each phenomenon (Table 1). We show in Fig.2 a brief sketch on the history of QCD and typical QCD effective models [1]. In ’80s, chiral symmetry breaking was the central issue. The chiral bag model, the NJL model and the σ model were formulated with referring chiral symmetry. In ’90s, on the other hand, the confinement physics is providing an important current of the hadron physics, since recent lattice QCD studies [4] shed a light on the confinement mechanism: the origin of color confinement can be recognized as the dual Higgs mechanism by monopole condensation, and the nonperturbative QCD vacuum is regarded as the dual superconductor [5]. The dual Ginzburg-Landau theory [6-12] was formulated from QCD with this picture, and provides a useful framework for the study of the nonperturbative QCD.

2. Color Confinement and Dual Higgs Mechanism

We briefly show the modern current of the confinement physics. The QCD vacuum can be regarded as the dual version of the superconductor, which was firstly pointed out by Nambu, ’t Hooft and Mandelstam [13]. Here, the “dual version” means the interchange between the electric and magnetic sectors. With referring Table 2 and Fig.3, we compare the ordinary electromagnetic system, the superconductor and the nonperturbative QCD vacuum regarded as the dual superconductor.

In the ordinary electromagnetism in the Coulomb phase, both electric flux and magnetic flux are conserved, respectively. The electric-flux conservation is guarantied by the ordinary gauge symmetry. On the other hand, the magnetic flux conservation is originated from the dual gauge symmetry [7-12], which

is the generalized version of the Bianchi identity. As for the inter-charge potential in the Coulomb phase, both electric and magnetic potentials are Coulomb-type.

The superconductor in the Higgs phase is characterized by electric-charge condensation, which leads to the Higgs mechanism or spontaneous breaking of the ordinary gauge symmetry, and therefore the electric flux is no more conserved. In such a system obeying the London equation, the electric inter-charge potential becomes short-range Yukawa-type similarly in the electro-weak unified theory. On the other hand, the dual gauge symmetry is not broken, so that the magnetic flux is conserved, but is squeezed like a one-dimensional flux tube due to the Meissner effect. As the result, the magnetic inter-charge potential becomes linearly rising like a condenser.

The nonperturbative QCD vacuum regarded as the dual Higgs phase is characterized by color-magnetic monopole condensation, which leads to the spontaneous breaking of the dual gauge symmetry. Therefore, color-magnetic flux is not conserved, and the magnetic inter-charge potential becomes short-range Yukawa-type. Note that the ordinary gauge symmetry is not broken by such monopole condensation. Therefore, color-electric flux is conserved, but is squeezed like a one-dimensional flux-tube or a string as a result of the dual Meissner effect. Thus, the hadron flux-tube is formed in the monopole-condensed QCD vacuum, and the electric inter-charge potential becomes linearly rising, which confines the color-electric charges [7-11].

As a remarkable fact in the duality physics, these are two “see-saw relations” between the electric and magnetic sectors.

- (1) There appears the Dirac condition $eg = 4\pi$ [7] in QCD. Here, unit electric charge e is the gauge coupling constant, and unit magnetic charge g is the dual gauge coupling constant. Therefore, a strong-coupling system in one sector corresponds to a weak-coupling system in the other sector.
- (2) As shown in Fig.3, the long-range confinement system in one sector corresponds to a short-range (Yukawa-type) interaction system in the other sector.

Let us consider usefulness of the latter “see-saw relation”. One faces highly non-local properties among the color-electric charges in the QCD vacuum because of the long-range linear confinement potential. Then, the direct formulation among the electric-charged variables would be difficult due to the non-locality, which seems to be a destiny in the long-distance confinement physics. However, one finds a short-range Yukawa potential in the magnetic sector, and therefore the electric-confinement system can be approximated by a local formulation among magnetic-charged variables. Thus, the confinement system, which seems highly non-local, can be described by a short-range interaction theory using the dual variables. This is the most attractive point in the dual Higgs theory.

Color-magnetic monopole condensation is necessary for color confinement in the dual Higgs theory. As for the appearance of color-magnetic monopoles in QCD, 't Hooft proposed an interesting idea of the abelian gauge fixing [14,15], which is defined by the diagonalization of a gauge-dependent variable. In this gauge, QCD is reduced into an abelian gauge theory with the color-magnetic monopole [7,14], which will be called as QCD-monopoles hereafter. Similar to the 't Hooft-Polyakov monopole [3] in the Grand Unified theory (GUT), the QCD-monopole appears from a hedgehog configuration corresponding to the non-trivial homotopy group $\pi_2(\text{SU}(N_c)/\text{U}(1)^{N_c-1}) = \mathbb{Z}^{N_c-1}$ on the nonabelian manifold.

Many recent studies based on the lattice QCD with 't Hooft abelian gauge fixing [4,11,16-18] show QCD-monopole condensation in the confinement phase, and strongly support abelian dominance and monopole dominance for the nonperturbative QCD (NP-QCD), e.g., linear confinement potential, dynamical chiral symmetry breaking (D χ SB) and instantons. Here, abelian dominance means that QCD

phenomena is described only by abelian variables in the abelian gauge. Monopole dominance is more strict, and means that the essence of NP-QCD is described only by the singular monopole part of abelian variables [4,11,16-18].

Figure 4 is the schematic explanation on abelian dominance and monopole dominance observed in the lattice QCD.

- (a) Without gauge fixing, it is very difficult to extract relevant degrees of freedom in NP-QCD.
- (b) In the abelian gauge, only U(1) gauge degrees of freedom including monopole is relevant for NP-QCD: abelian dominance. Here, monopole condensation is observed as the appearance of a very long monopole loop in the confinement phase. On the other hand, off-diagonal parts does not contribute to NP-QCD.
- (c) The U(1)-variable can be separated into the regular photon part and the singular monopole part [19]. The monopole part leads to NP-QCD (confinement, D χ SB, instanton): monopole dominance. On the other hand, the photon part is almost trivial similar to the QED system.

Thus, the condensed monopole in the 't Hooft abelian gauge is nothing but the relevant collective mode for NP-QCD, and therefore the NP-QCD vacuum can be identified as the dual-superconductor in a realistic sense.

3. Asymptotic Freedom in the Dual Ginzburg-Landau Theory

3-1. Dual Ginzburg-Landau Theory

The dual Ginzburg-Landau (DGL) theory is the QCD effective theory based on the dual Higgs mechanism, and can be derived from the QCD lagrangian with considering the QCD nature: monopole condensation and abelian dominance [4]. The DGL lagrangian [11,12] for the pure-gauge system is described with the dual gauge field B_μ and the QCD-monopole field χ ,

$$\mathcal{L}_{\text{DGL}} = \text{tr} \left\{ -\frac{1}{2}(\partial_\mu B_\nu - \partial_\nu B_\mu)^2 + [\hat{D}_\mu, \chi]^\dagger [\hat{D}^\mu, \chi] - \lambda(\chi^\dagger \chi - v^2)^2 \right\}, \quad (2)$$

where $\hat{D}_\mu \equiv \partial_\mu + igB_\mu$ is the dual covariant derivative including the dual gauge coupling constant $g = 4\pi/e$.

The dual gauge field $B_\mu \equiv \vec{B}_\mu \cdot \vec{H} = B_\mu^3 T^3 + B_\mu^8 T^8$ is defined on the dual gauge manifold $U(1)_m^3 \times U(1)_m^8$ [11,12], which is the dual space of the maximal torus subgroup $U(1)_e^3 \times U(1)_e^8$ embedded in the original gauge group $SU(3)_c$. The abelian field strength tensor is written as $F_{\mu\nu} = *(\partial \wedge B)_{\mu\nu}$, so that the role of the electric and magnetic fields are interchanged in comparison with the ordinary A_μ description.

The QCD-monopole field χ is defined as $\chi \equiv \sqrt{2} \sum_{\alpha=1}^3 \chi_\alpha E_\alpha$ [11,12] using $E_1 \equiv \frac{1}{\sqrt{2}}(T_6 + iT_7)$, $E_2 \equiv \frac{1}{\sqrt{2}}(T_4 - iT_5)$ and $E_3 \equiv \frac{1}{\sqrt{2}}(T_1 + iT_2)$. Here, χ_α has the magnetic charge $g\vec{\alpha}$ proportional to the root vector $\vec{\alpha}$. In the QCD-monopole condensed vacuum with $|\chi_\alpha| = v$, the dual gauge symmetry $U(1)_m^3 \times U(1)_m^8$ is spontaneously broken. Through the dual Higgs mechanism, the dual gauge field B_μ acquires its mass $m_B = \sqrt{3}gv$, whose inverse provides the radius of the hadron flux tube [7], and the dual Meissner effect causes the color-electric field excluded from the QCD vacuum, which leads to color confinement. The QCD-monopole fluctuations $\tilde{\chi}_\alpha \equiv \chi_\alpha - v$ ($\alpha=1,2,3$) also acquire their mass $m_\chi = 2\sqrt{\lambda}v$ in the QCD-monopole condensed vacuum. As a relevant prediction, only one QCD-monopole fluctuation $\tilde{\chi} \equiv \sum_{\alpha=1}^3 \tilde{\chi}_\alpha$ appears as a color-singlet scalar glueball in the confinement phase, although the dual gauge field B_μ and the other two combinations of the QCD-monopole fluctuation are not color-singlet and cannot be observed [11,17,18].

The DGL theory reproduces confinement properties like the inter-quark potential and the hadron flux-tube formation. We studied effects of the flux-tube breaking by the light quark-pair creation in the DGL theory, and derived the infrared screened inter-quark potential [7-9], which is observed in the lattice QCD with dynamical quarks. We studied also the dynamical chiral-symmetry breaking (D χ SB), which is also an important feature in the nonperturbative QCD, by solving the Schwinger-Dyson equation for the dynamical quark [7-9]. The quark-mass generation is brought by QCD-monopole condensation, which suggests the close relation between D χ SB and color confinement. Thus, the DGL theory provides not only the confinement properties but also D χ SB and its related quantities like the constituent quark mass, the chiral condensate and the pion decay constant [7-11,20].

3-2. Asymptotic Behavior in DGL Theory

In the abelian gauge fixing, the Dirac condition $eg = 4\pi$ for the dual gauge coupling constant g is derived [14,7] as a geometrical constraint on the nonabelian manifold using the similar argument on the GUT-monopole [3]. Since the Dirac condition is a geometrical relation, it is renormalization group invariant. As the first category of the “see-saw relation” in Section 2, g becomes small for the strong coupling region with a large coupling constant e [7,8].

The DGL theory in the pure gauge is renormalizable similar to the abelian Higgs model [2], and is not asymptotically free on g in view of the renormalization group: $g(p^2)$ is increasing function of p^2 . Hence, asymptotic freedom is expected for the QCD gauge coupling constant e owing to the Dirac condition: $e(p^2)$ defined by $e(p^2)g(p^2) = 4\pi$ is decreasing function of p^2 . Thus, the DGL theory qualitatively shows asymptotic freedom on the QCD gauge coupling e [7,8]. This asymptotic behavior in the DGL theory is consistent with QCD qualitatively, and seems a desirable feature for an effective theory of QCD.

Next, we attempt to calculate the β -function and the running coupling constant from the polarization tensor $\Pi_{\mu\nu}^{ab}(p)$ of the dual gauge field B_μ in the DGL theory. In particular, we are interested in the infrared region ($p \lesssim 1\text{GeV}$), where the perturbative QCD calculation is not reliable.

With the dimensional regularization [2,3] by shifting the space-time dimension as $d = 4 - \epsilon$, we calculate the simplest nontrivial radiative correction from the QCD-monopole loop diagrams as shown in Fig.5,

$$\Pi_{\mu\nu}^{ab}(p) = -\frac{1}{32\pi^2}\delta^{ab}(g\mu^{-\epsilon})^2(p^2 g_{\mu\nu} - p_\mu p_\nu)\frac{1}{\epsilon} + O(\epsilon^0), \quad (3)$$

where g is the bare dual-gauge coupling and μ the renormalization point. In the minimum subtraction scheme [3], the $O(1/\epsilon)$ divergence is eliminated by the counter term contribution,

$$\Pi_{\mu\nu}^{Cab}(p) = -(Z_3 - 1)\delta^{ab}(p^2 g_{\mu\nu} - p_\mu p_\nu), \quad (4)$$

where Z_3 is the wave-function renormalization constant [2] of the dual gauge field B_μ ,

$$Z_3(\mu) = 1 - \frac{(g\mu^{-\epsilon})^2}{32\pi^2}\frac{1}{\epsilon}. \quad (5)$$

Because of the Ward identity $Z_1 = Z_2$ [2], the renormalized coupling constant is simply given by $g(\mu) = Z_3(\mu)^{1/2}g$. The β -function [2,3] is then expressed as

$$\beta \equiv \mu \frac{d}{d\mu} g(\mu) = \mu \frac{d}{d\mu} \{Z_3(\mu)^{1/2}g\} = \frac{1}{2\epsilon} g(\mu)^3 + O[g(\mu)^5], \quad (6)$$

which determines the behavior of the running coupling $g(\mu)$ as

$$\frac{1}{g^2(\mu)} = \frac{1}{g^2(\mu_0)} - \frac{1}{32\pi^2} \ln(\mu^2/\mu_0^2) \quad (7)$$

within the leading order.

By summation of the multi-polarization diagrams, one obtains the final formula for the running coupling $g(\mu)$ including the higher order correction,

$$g^2(\mu) = g^2(\mu_0) - \frac{1}{32\pi^2} \ln(\mu^2/\mu_0^2). \quad (8)$$

In the DGL theory, the QCD gauge coupling $e(\mu)$ defined by $e(\mu)g(\mu) = 4\pi$ is simply expressed as

$$e^2(\mu) = e^2(\mu_0) \left\{ 1 + \frac{1}{e^2(\mu_0)} \ln(\mu/\mu_0) \right\}^{-1}. \quad (9)$$

We show in Fig.6 the running coupling constants, $g(\mu)$ and $e(\mu)$, in the DGL theory with the parameter set in Ref.[9]: $m_\chi=1.67\text{GeV}$. Similarly in QED or the abelian Higgs model, we have imposed the renormalization condition as $g(\mu=2m_\chi)=7.9$ [$e(\mu=2m_\chi)=1.59$], which is taken to be consistent with the parameter $g=6.28$ ($e=2.0$) in Ref.[9] in the infrared region (see Fig.6).

As a model prediction in the DGL theory, the gauge coupling $e(\mu)$ behaves as “walking coupling constant”, which means the slowly varying running coupling, even in the infrared region as $\mu \lesssim 1\text{GeV}$.

4. Correlation between Instanton and QCD-monopole

4-1. Analytical Study

The instanton is another important topological object in the nonabelian gauge theory; $\pi_3(\text{SU}(N_c)) = \mathbb{Z}_\infty$ [11,17,18,21]. Recent lattice studies [4,11,16-18] indicate abelian dominance for the nonperturbative quantities in the maximally abelian (MA) gauge and in the Polyakov gauge. If the system is completely described only by the abelian field, the instanton would lose the topological basis for its existence, and therefore it seems unable to survive in the abelian manifold. However, even in the abelian gauge, nonabelian components remain relatively large around the QCD-monopoles, which are nothing but the topological defects, so that instantons are expected to survive only around the QCD-monopole trajectories in the abelian-dominant system [11,17,18].

We have pointed out such a close relation between instantons and QCD-monopoles in the continuum $\text{SU}(2)$ gauge theory in the Polyakov-like gauge, where $A_4(x)$ is diagonalized [8,17,18]. Using the 't Hooft symbol $\bar{\eta}^{a\mu\nu}$, the multi-instanton solution is written as [11,17,18] $A^\mu(x) = i\bar{\eta}^{a\mu\nu} \frac{\tau^a}{2} \partial^\nu \ln \left(1 + \sum_k \frac{a_k^2}{|x-x_k|^2} \right)$, where $x_k^\mu \equiv (\mathbf{x}_k, t_k)$ and a_k denote the center coordinate and the size of k -th instanton, respectively. Near the center of k -th instanton, $A_4(x)$ takes a hedgehog configuration, $A_4(x) \simeq i \frac{\tau^a (\mathbf{x}-\mathbf{x}_k)^a}{|x-x_k|^2}$. In the Polyakov-like gauge, $A_4(x)$ is diagonalized by a singular gauge transformation [7,18], which leads to the QCD-monopole trajectory on $A_4(x) = 0$: $\mathbf{x} \simeq \mathbf{x}_k$. Hence, the QCD-monopole trajectory penetrates each instanton center along the temporal direction. In other words, instantons only live along the QCD-monopole trajectory.

In the multi-instanton system [8,17,18,22,23], the QCD-monopole trajectory tends to be highly complicated and unstable against a small fluctuation of the location or the size of instantons [17,18] even at the classical level. The QCD-monopole trajectory has loops or a folded structure, and the topology of the trajectory is often changed due to a small fluctuation of instantons [17,18].

We show in Fig.7 an example of the QCD-monopole trajectory in the multi-instanton system, where all instantons are put on the zt -plane for simplicity. The contour denotes the magnitude of the topological density. Each instanton attaches the QCD-monopole trajectory. As a remarkable feature in the Polyakov-like gauge, the monopole favors the high topological density region, “mountain”. Each monopole trajectory walks crossing tops of the mountain. On the other hand, anti-monopole with the opposite color-magnetic charge favors the low topological density region, “valley”. Thus, the strong local correlation is found between the instanton and the QCD-monopole.

Furthermore, quantum fluctuation should make the QCD-monopole world line more complicated and more unstable, which leads to appearance of a long complicated trajectory as a result. Thus, instantons would contribute to promote monopole condensation, which is signaled by a long complicated monopole loop as shown in Fig.8 in the lattice QCD simulation [4]

4-2. Lattice Study

We study the correlation between instantons and QCD-monopoles [16,17,18,24] in the maximally abelian (MA) gauge and in the Polyakov gauge using the SU(2) lattice with 16^4 and $\beta = 2.4$. The SU(2) link variable can be separated into the singular monopole part and the regular photon part [16,17,18]. We find that instantons and anti-instantons exist only in the monopole sector in the abelian gauges, which means monopole dominance for the topological charge [16,17,18]. Monopole dominance for the $U_A(1)$ anomaly is also expected.

We study also the finite-temperature system using the $16^3 \times 4$ lattice with various β around $\beta_c \simeq 2.3$ [25]. Monopole dominance for the instanton is found in the finite-temperature confinement phase. Near the critical temperature $\beta_c \simeq 2.3$, we observe a large reduction of the number of instantons and anti-instantons. In the deconfinement phase ($\beta > \beta_c$), instantons vanish as well as QCD-monopole condensation. Hence, instantons are expected to survive only around the condensed QCD-monopole trajectories [25] (see Fig.10).

Finally, we study the correlation between instanton number and QCD-monopole loop length. There appears a very long QCD-monopole loop in each lattice gauge configuration in the confinement phase as shown in Fig.8. Hence, a very long QCD-monopole loop is regarded as a signal of the confinement in the lattice QCD [4]. As shown in Fig.9, a linear correlation is found between the total monopole-loop length L and $I_Q \equiv \int d^4x |\text{tr}(G_{\mu\nu}\tilde{G}_{\mu\nu})|$, which corresponds to the total number N_{tot} of instantons and anti-instantons.

From the above results, we propose the following conjecture. Each instanton accompanies a small monopole loop nearby [22,24], whose length would be proportional to the instanton size. When N_{tot} is large enough, these monopole loops overlap, and there appears a very long QCD-monopole trajectory, which bonds neighboring instantons as shown in Fig.10. Such a monopole clustering leads to monopole condensation and color confinement. Thus, multi-instanton is expected to provide a source of the color confinement through the monopole clustering [17,18].

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